

Photoabsorption on nuclei*

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Abstract

We calculate the total photoabsorption cross section on nuclei in the energy range from 300 MeV to 1 GeV within the framework of a semi-classical phase space model. Besides medium modifications like Fermi motion and Pauli blocking we focus on the collision broadening of the involved resonances. The resonance contributions to the elementary cross section are fixed by fits to partial wave amplitudes of pion photoproduction. The cross sections for $NR \rightarrow NN$, needed for the calculation of collision broadening, are obtained by detailed balance from a fit to $NN \rightarrow NN\pi$ cross sections. We show that a reasonable collision broadening is not able to explain the experimentally observed disappearance of the $D_{13}(1520)$ -resonance in the photoabsorption cross section on nuclei.

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1 Introduction

The total nuclear photoabsorption cross section in the first and second nucleon resonance region which was recently measured in Mainz [1] and Frascati

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[2, 3] shows clear medium modifications compared to the elementary cross section on the proton and deuteron [4, 5, 6]. One observes a strong broadening of the Δ -peak and the disappearance of the higher resonances D_{13} and F_{15} while the total cross section per nucleon is almost independent of the mass of the nucleus.

In the region of the Δ -resonance there are calculations within the Δ -hole model [7, 8] that are able to describe the experimental data up to 500 MeV. For higher energies there have been attempts [3, 9, 10] to deduce the collision broadening of the resonances from the experimental data by fitting the total photonuclear cross section; these studies lead to significantly increased resonance widths in nuclei.

We present now a consistent calculation of the photon-nucleus reaction over the whole energy range from 300 MeV to 1 GeV within the framework of a semi-classical phase space model. The microscopic input to our model calculation is identical to that used in the BUU transport model [11, 12] which has been very successfully applied to the description of heavy-ion collisions up to bombarding energies of 2 GeV/A [13], pion-nucleus reactions [14] and also photoproduction of pions and etas [15]. Our calculation is based on the assumption that the total photonuclear cross section is the incoherent sum of contributions from all nucleons where we neglect possible shadowing effects. Besides more or less trivial medium modifications like Fermi motion and Pauli blocking we investigate the effect of collision broadening for the involved resonances. This may lead to a better understanding of the behaviour of nucleon resonances in nuclear matter.

In section 2 we start with the presentation of the cross sections that we use for the interaction of resonances with nucleons. The results for the collision widths induced by these cross sections are shown in section 3. The elementary photoabsorption cross section on the nucleon is discussed in section 4. Finally the results for the total photoabsorption cross section on nuclei are presented in section 5.

2 Resonance-nucleon interactions

We use here resonance parameterizations that we have recently introduced for use in the collision term for BUU-calculations that includes baryonic resonances up to a mass of about 2 GeV. The used decay widths and cross

sections are described in detail in [13]. The one-pion decay width is parameterized according to [12]:

$$\Gamma(q) = \Gamma_0 \left(\frac{q}{q_r} \right)^{2l+1} \left(\frac{q_r^2 + c^2}{q^2 + c^2} \right)^{l+1} \quad (1)$$

where

$$c^2 = (M_R - M_N - m_\pi)^2 + \frac{\Gamma_0^2}{4} \quad . \quad (2)$$

Here, q and q_r denote the cms pion momenta of resonances with mass M and M_R respectively where M_R is the mass at the pole of the resonance. l is the angular momentum of the πN -system, M_N the nucleon mass and m_π the pion mass. Γ_0 is the decay width at the pole of the resonance. The used resonance parameters can be found in [13].

For the Δ -resonance we use the Moniz-parameterization [8]:

$$\Gamma_\Delta(q) = \Gamma_0 \frac{M_\Delta}{M} \left(\frac{q}{q_r} \right)^3 \left(\frac{q_r^2 + c^2}{q^2 + c^2} \right)^2 \quad \text{with } c = 0.3 \text{ GeV} \quad , \quad (3)$$

because this parameterization gives a good description of the P_{33} -multipole of elastic πN -scattering within a Breit-Wigner approximation [16].

The two-pion decay is parameterized as a two-step process:

$$R \rightarrow r a \rightarrow N \pi \pi$$

where r stands for a baryonic or mesonic resonance and a is a pion or nucleon. The two-pion decay width is then given by a phase space weighted integral over the mass distribution of the intermediate resonance:

$$\begin{aligned} \Gamma_{R \rightarrow r a}(M) &= \frac{k}{M} \int_0^{M-m_a} d\mu p_f \frac{2}{\pi} \frac{\mu^2 \Gamma_{r,tot}(\mu)}{(\mu^2 - m_r^2)^2 + \mu^2 \Gamma_{r,tot}^2(\mu)} \\ &\times \frac{(M_R - M_N - 2m_\pi)^2 + c^2}{(M - M_N - 2m_\pi)^2 + c^2} \quad \text{with } c = 0.3 \text{ GeV} \quad , \quad (4) \end{aligned}$$

where p_f denotes the cms momentum of a and r and μ the mass of the intermediate resonance. The constant k follows from the knowledge of $\Gamma_{R \rightarrow r a}(M_R)$. The parameters of the mesons and the parameterizations of their decay widths can be found in [13].

The parameterizations of the partial widths of the $N(1535)$ were fitted to the eta photoproduction cross section on the proton and are described in [17].

For the calculation of collision broadening one needs to know the cross sections for interactions of resonances with nucleons. The cross sections for resonance absorption processes ($N R \rightarrow N N$) are obtained by detailed balance from one-pion production in nucleon-nucleon collisions under the assumption that the main contribution to $N N \rightarrow N N \pi$ comes from two-step processes $N N \rightarrow N R \rightarrow N N \pi$. The contribution of the $\Delta(1232)$ -resonance is adopted from [18] where we have only replaced the parameterization of the Δ -width (equation (3)). This cross section for $N N \rightarrow N \Delta(1232)$ describes experimental massdifferential cross sections reasonably well [18, 13].

For the higher resonances massdifferential data are not available. Therefore we make the ansatz that the matrix element for $N N \rightarrow N R$ is constant and equal for all higher resonances. This matrix element is obtained by fitting total $N N \rightarrow N N \pi$ cross sections. In order to get a better description of the threshold behaviour a background term is added. From figure 1 one sees that the contributions coming from the higher resonances are small compared to that from the $\Delta(1232)$.

For the resonance absorption process $N R \rightarrow N N$ we obtain the cross section:

$$\sigma_{NR \rightarrow NN} = S I \frac{2}{2J_R + 1} \frac{p_f |\mathcal{M}_{NN \rightarrow NR}|^2}{16\pi p_i s}, \quad (5)$$

where J_R denotes the spin of the resonance, p_i and p_f are the cms momenta of the incoming and outgoing particles respectively and \sqrt{s} is the total energy in the cms. The isospin factors I result from an incoherent sum over all contributing isospin channels under the assumption that the absolute value of the matrix element $\mathcal{M}_{NN \rightarrow NR}$ does not depend on the total isospin. They are given in table 1. The factor S reads:

$$S = \begin{cases} \frac{1}{2} & , \text{ identical particles in the final state} \\ 1 & , \text{ else.} \end{cases}$$

The numerical value for the matrix element is fitted to the cross sections for $N N \rightarrow N N \pi$ resulting in:

$$\frac{|\mathcal{M}_{NN \rightarrow NR}|^2}{16\pi} = 18 \text{ mb GeV}^2. \quad (6)$$

		I
$N^+ \quad N^+$	$N^+ \quad N^+$	1
$N^+ \quad N^0$	$N^+ \quad N^0$	1/2
$N^+ \quad N^+$	$N^0 \quad \Delta^{++}$	3/4
$N^+ \quad N^+$	$N^+ \quad \Delta^+$	1/4
$N^+ \quad N^0$	$N^+ \quad \Delta^0$	1/4
$N^+ \quad \Delta^{++}$	$N^+ \quad \Delta^{++}$	1
$N^+ \quad \Delta^+$	$N^0 \quad \Delta^{++}$	3/8
$N^+ \quad \Delta^+$	$N^+ \quad \Delta^+$	5/8
$N^+ \quad \Delta^0$	$N^+ \quad \Delta^0$	1/2
$N^+ \quad \Delta^-$	$N^+ \quad \Delta^-$	5/8

Table 1: Isospin coefficients for baryon-baryon collisions. N stands for any isospin- $\frac{1}{2}$ particle, Δ for any isospin- $\frac{3}{2}$ one.

The matrix elements for the $N(1535)$ -resonance are fitted to $NN \rightarrow NN\eta$ cross sections. This gives [13]:

$$\frac{|\mathcal{M}_{pp \rightarrow pN^+(1535)}|^2}{16\pi} = 8 \text{ mb GeV}^2 \quad (7)$$

$$\text{and } \frac{|\mathcal{M}_{pn \rightarrow nN^+(1535)}|^2}{16\pi} = 20 \text{ mb GeV}^2 \quad . \quad (8)$$

The cross section for $N\Delta(1232) \rightarrow N\Delta(1232)$ is estimated by the calculation of the diagram shown in figure 2 (a) alone, because the process in figure 2 (b) with an intermediate on-shell pion is already included in our transport model via the two-step process. In figure 3 the resulting cross section for the one-step process is depicted together with the cross section for $\Delta(1232)N \rightarrow NN$.

The matrix elements for $NR \rightarrow NR$ (the same resonance in in- and outgoing channel) are estimated by using a parameterization of the elastic nucleon-nucleon scattering cross section [19]:

$$\sigma_{NN \rightarrow NN} = \left(\frac{35}{1 + \frac{\sqrt{s} - 2M_N}{\text{GeV}}} + 20 \right) \text{ mb} \quad . \quad (9)$$

The assumption of an isotropic angular dependence gives the corresponding squared matrix element:

$$|\mathcal{M}_{NN \rightarrow NN}|^2 = 16 \pi s \left(\frac{35}{1 + \frac{\sqrt{s-2M_N}}{\text{GeV}}} + 20 \right) \text{ mb} \quad . \quad (10)$$

This matrix element is now used for the $NR \rightarrow NR$ cross section:

$$\sigma_{NR \rightarrow NR} = I \frac{|\mathcal{M}_{NN \rightarrow NN}|^2}{16 \pi p_i s} \int d\mu p_f \frac{2}{\pi} \frac{\mu^2 \Gamma_R(\mu)}{(\mu^2 - M_R^2)^2 + \mu^2 \Gamma_R^2(\mu)} \quad , \quad (11)$$

where I is again the isospin coefficient from table 1 and μ denotes the mass of the outgoing resonance. For processes with a change of the resonance we use the same matrix element as for $NR \rightarrow NN$. This gives:

$$\sigma_{NR \rightarrow NR'} = I \frac{2}{2J_R + 1} \frac{|\mathcal{M}_{NN \rightarrow NR}|^2}{16 \pi p_i s} \int d\mu p_f \frac{2}{\pi} \frac{\mu^2 \Gamma_{R'}(\mu)}{(\mu^2 - M_{R'}^2)^2 + \mu^2 \Gamma_{R'}^2(\mu)} \quad . \quad (12)$$

3 Collision broadening

The basic concepts of collision broadening of resonances in nuclei are described in [10]. We have used the cross sections for the interaction of resonances with nucleons which were described in the preceding section to calculate the collision widths of the resonances that are important for photonuclear reactions. These are the $P_{33}(1232)$, the $D_{13}(1520)$ and the $F_{15}(1680)$. The $S_{11}(1535)$ is important for the calculation of etaproduction.

Since nucleon final states coming from spontaneous decays of resonances can be Pauli blocked in the nuclear medium there is also a reduction of the width. The total in-medium width is therefore:

$$\Gamma_{tot}^{med} = \Gamma_{spon}^{med} + \Gamma_{coll}^{med} \quad , \quad (13)$$

where Γ_{spon}^{med} stands for the sum of the one-pion-, two-pion- and eta-width.

The in-medium widths are calculated in nuclear matter and then applied to finite nuclei by means of a local density approximation. The collision

width coming from the process $N R \rightarrow N N$ for a resonance with mass M_R and momentum p_R at nucleon density ρ is thus given by:

$$\Gamma_{N_1 R \rightarrow N_2 N_3}(M_R, p_R, \rho) = 4 \gamma \int_0^{p_F} \frac{d^3 p_{N1}}{(2\pi)^3} v_r \int d\Omega \frac{d\sigma_{NR \rightarrow NN}}{d\Omega} P_{N2} P_{N3} S, \quad (14)$$

where the relative velocity v_r is

$$v_r = \sqrt{\frac{(\mathbf{p}_R \cdot \mathbf{p}_{N1})^2 - M_R^2 M_N^2}{(\mathbf{p}_R \cdot \mathbf{p}_{N1})^2}}, \quad (15)$$

with \mathbf{p}_R and \mathbf{p}_{N1} being the 4-momenta of the resonance and incoming nucleon respectively. The Pauli blocking function

$$P_N = \Theta(|\vec{p}_N| - p_F) \quad (16)$$

takes into account that the momenta of the outgoing nucleons have to be outside of the Fermi sphere. A shadowing function S cuts off the in-medium cross section at high values:

$$S = \begin{cases} 1 & , \quad \sigma_{NR,tot} \leq \sigma_{max} \\ \frac{\sigma_{max}}{\sigma_{NR,tot}} & , \quad \sigma_{NR,tot} > \sigma_{max} \end{cases}, \quad (17)$$

where σ_{max} is

$$\sigma_{max} = \left(\frac{\rho_0}{\rho}\right)^{\frac{2}{3}} 80.4 \text{ mb} \quad . \quad (18)$$

Here,

$$\sqrt{\frac{80.4 \text{ mb}}{\pi}} = 1.6 \text{ fm}$$

is approximately equal to the spatial distance of neighboring nucleons at nucleon density $\rho_0 = 0.168/\text{fm}^3$. The results reported later are insensitive to an increase of σ_{max} .

The collision width for $N R \rightarrow N R'$ scattering is calculated analogously to equation (14) with an additional integration over the mass distribution of the outgoing resonance since the momentum of the outgoing nucleon entering the Pauli blocking function depends on the mass μ of the outgoing resonance:

$$\begin{aligned} \Gamma_{N_1 R \rightarrow N_2 R'}(M_R, p_R, \rho) &= 4 \gamma \int_0^{p_F} \frac{d^3 p_{N1}}{(2\pi)^3} v_r \int d\Omega \\ &\times \int d\mu \sum_{R'} \frac{d^2 \sigma_{NR \rightarrow NR'}}{d\Omega d\mu} P_{N2} S \quad . \end{aligned} \quad (19)$$

We neglect the selfconsistency in this equation due to the width dependence of the cross sections on the rhs because the effect of the in-medium width in equation (14) compared to the vacuum width is very small. Moreover one has to keep in mind that the uncertainties in the cross sections for $N R \rightarrow N R'$ are very large.

The in-medium one-pion decay width is obtained by averaging over the decay angle:

$$\Gamma_{R \rightarrow N\pi}(M_R, p_R, \rho) = \frac{1}{2} \int_{-1}^1 d\cos(\theta) \Gamma_{0,N\pi}(M_R) P_N \quad , \quad (20)$$

where we assume an isotropic decay in the rest frame of the resonance. θ is the angle between outgoing nucleon momentum and boost axis in the rest frame of the resonance and $\Gamma_{0,N\pi}$ the vacuum one-pion decay width. The in-medium eta width is calculated analogously.

The two-pion decay width is parameterized as a two-step process with a first decay into a baryonic or mesonic resonance and a pion or nucleon (section 2). For the decay into a baryonic resonance and a pion we assume no medium modifications while the decay into a nucleon and a mesonic resonance r is modified due to possible Pauli blocking of the nucleon final state:

$$\Gamma_{R \rightarrow r N}(M_R, p_R, \rho) = \frac{1}{2} \int_{-1}^1 d\cos(\theta) \int dm_r \frac{d\Gamma_{0,Nr}(M_R)}{dm_r} P_N \quad , \quad (21)$$

where m_r denotes the mass of the outgoing mesonic resonance.

In figure 4 we show the different contributions to the in-medium width of the $\Delta(1232)$ at nucleon densities ρ_0 and $\rho_0/2$ in isospin symmetric nuclear matter. Here the momentum of the resonance is related to its mass by the requirement that the resonance was created by photoabsorption on a free nucleon at rest. For comparison the vacuum width is also shown.

At ρ_0 the collision width coming from resonance absorption $N \Delta \rightarrow N N$ is at the pole of the resonance ($M=1.232$ GeV) about 25 MeV. This partial width decreases with increasing mass. The contribution from $N \Delta \rightarrow N \Delta$ has - averaged over the mass distribution - about the same size as the one from $N \Delta \rightarrow N N$. However, here we get a strong increase of $\Gamma_{N\Delta \rightarrow N\Delta}$ with increasing mass because of the phase space weighted integral over the mass distribution of the outgoing Δ -resonance (equation (19)).

Our calculations of the collision widths for the Δ -resonance are in agreement with the estimates given by Kondratyuk et al. (chapter 4 in [10]). The

process $N \Delta \rightarrow N R$ is negligible because the energy of the $N \Delta$ -system is too low.

In the region of the resonance pole the total in-medium width is almost independent of the nucleon density since collision broadening and Pauli reduction of the free width nearly compensate. Thus, at the resonance pole the net broadening compared to the vacuum width is very small; at about 100 MeV above the pole the width has grown by about 50 MeV, mainly due to the $N \Delta \rightarrow N \Delta$ scattering process.

In figure 5 the in-medium widths of the $N(1520)$, the $N(1535)$ and the $N(1680)$ are compared with the vacuum widths and split up into their partial widths. The very strong rise of the width of the $N(1520)$ is due to the opening of the ρ -channel which is probably not cut off fast enough. Nevertheless, in all cases the collision widths at the poles of the resonances are only of the order 20 - 40 MeV which leads to a small net broadening because the Pauli blocking of the free width is less important than in the case of the Δ -resonance. Notice that the collision widths of those higher lying resonances increase less significantly with mass than those of the Δ .

The main contribution to the collision widths comes from the process $N R \rightarrow N R$ for which we estimated the cross section by adopting the matrix element from elastic nucleon nucleon scattering. The partial widths of $N R \rightarrow N N$ are very small (< 10 MeV) as to be expected from figure 1. Even if one dropped the assumption that the matrix element is equal for all higher resonances and set the matrix element of the $D_{13}(1520)$ to the maximum possible value which is in line with the $N N \rightarrow N N \pi$ data, one would only get an additional broadening of the $D_{13}(1520)$ of about 30 MeV. Therefore a collision width of 300 MeV for the $D_{13}(1520)$ as given in [9, 10] seems to be far from being realistic.

In order to check the validity of the local density approximation used for the application of the discussed in-medium widths we also performed a calculation of the Δ -width in finite nuclei. Here we defined the density to be the density at the location of creation of the resonance. It turned out that there was a difference between this calculation and the one in nuclear matter only for very low mass Δ 's because the width of these Δ 's is small enough for them to travel through a relevant density gradient within their lifetimes. For Δ 's in the region of the resonance pole the mean free path is only about 0.5 fm leading to a sensible applicability of the local density approximation.

In summary, the effective collisional broadening of the resonances never

amounts to more than 10% at the resonance pole.

4 The total photoabsorption cross section on the nucleon

In [17] the decomposition of the total photoabsorption cross section on the nucleon into the different channels and the resonance contributions is described in detail. For the one-pion production cross section we use partial-wave amplitudes as given by [20] and fit the resonance contributions to these amplitudes because - especially in the region of the Δ -resonance - interference terms with the background are quite important. An incoherent decomposition of the total photoabsorption cross section into resonance and background contributions as done by Kondratyuk et al. [10] should not be used if one wants to investigate possible modifications of the resonance contributions in nuclei.

While the one-pion production cross sections can nicely be decomposed into Breit-Wigner type resonance contributions and a smooth background, the structure of the two-pion production cross sections is not described by the resonance contributions that are induced by the two-pion decay widths of the resonances [17]. The difference between the experimental cross section and the Breit-Wigner type resonance contributions is treated as background, where the momenta of the outgoing particles are distributed according to three-body phase space. Therefore the only medium modification is the possible Pauli blocking of the outgoing nucleon.

In figure 6 we show the sum of the contributions to the total photoabsorption cross section on the proton with the experimental data [5, 21]. The good agreement indicates that all important channels are included. For photon energies above 600 MeV the 'background' for $\gamma p \rightarrow N \pi \pi$ amounts for about one half of the total cross section.

For the neutron the two-pion background was fitted to the absorption cross section [17]. In the region of the Δ -resonance the sum of the one-pion cross sections [20] alone is about $150 \mu\text{b}$ larger than the total absorption cross section published by Armstrong et al. [6]. Since the photoabsorption cross section on the deuteron measured by Armstrong et al. has recently been confirmed by the DAPHNE collaboration [4] we assume this discrepancy to

arise from the extraction of the neutron cross section from the deuteron data.

Unfortunately there is yet no calculation of the photoabsorption on the deuteron in the energy range of interest. However, there is a recent calculation of quasifree pion photoproduction on the deuteron in the Δ -region within the spectator nucleon model [22]. Here, the cross section for $\gamma d \rightarrow N N \pi$ at the Δ -peak is about 1100 μb . This result is at variance with the measured total photoabsorption cross section on the deuteron being only about 900 μb [4, 6].

4.1 Medium modifications

We use the following medium modifications for the elementary photon-nucleon cross section:

- The vacuum width appearing in the resonance propagators is replaced by the in-medium width as calculated in section 3.
- The collision width gives a Breit-Wigner type contribution to the absorption cross section [17].
- For the Δ -resonance the difference between nucleon and Δ -potential causes a real part of the self energy Π to be used in the resonance propagator:

$$\text{Re } \Pi = 2 E_{\Delta} (U_N - U_{\Delta}) \quad . \quad (22)$$

- Nucleon final states can be Pauli blocked.

5 Results

We now combine the ingredients described in the preceding sections in a calculation of the total photoabsorption cross section on nuclei using the local Thomas-Fermi model with realistic density distributions. Neglecting shadowing effects the total cross section on a nucleus within our model then reads:

$$\sigma_{abs} = \int d^3r \int^{p_F} 4 \frac{d^3p_N}{(2\pi)^3} \frac{k_{\gamma N}}{k_{\gamma A}} \frac{M_N}{E_N} \left\{ \int \left[\frac{Z}{A} \left(\frac{d\sigma_{\gamma p \rightarrow N\pi}^{med}}{d\Omega} + \frac{d\sigma_{\gamma p \rightarrow p\eta}^{med}}{d\Omega} \right) \right. \right.$$

$$\begin{aligned}
& + \frac{A-Z}{A} \left(\frac{d\sigma_{\gamma n \rightarrow N\pi}^{med}}{d\Omega} + \frac{d\sigma_{\gamma n \rightarrow n\eta}^{med}}{d\Omega} \right) d\Omega \\
& + \int \left[\frac{Z}{A} \frac{d\sigma_{\gamma p \rightarrow N\pi\pi}^{med}}{d\Omega' dp'_N} + \frac{A-Z}{A} \frac{d\sigma_{\gamma n \rightarrow N\pi\pi}^{med}}{d\Omega' dp'_N} \right] d\Omega' dp'_N \Big\} \quad , \quad (23)
\end{aligned}$$

where $p_F[\rho(r)]$ denotes the local Fermi momentum [17], $k_{\gamma N}$ is the photon energy in the rest frame of the nucleon with momentum \vec{p}_N and $k_{\gamma A}$ the photon momentum in the rest frame of the nucleus. The factor

$$\frac{k_{\gamma N}}{k_{\gamma A}} \frac{M_N}{E_N}$$

results from the Lorentz transformation of the cross section [23].

Figure 7 shows the cross sections on ^{40}Ca that result from successive application of the medium modifications. Fermi motion alone leads to a damping of the Δ -peak by about $100 \mu\text{b}$ per nucleon. The structure in the region of the D_{13} -resonance is washed out but does not disappear. The peak of the F_{15} -resonance vanishes. Pauli blocking further decreases the Δ -peak by about $100 \mu\text{b}$ with the reduction of the cross section getting smaller at higher energies.

Using the in-medium widths for the resonances the Δ -peak is shifted to smaller energies and increased by about $70 \mu\text{b}$ per nucleon. The reason is given by the comparison of the Δ -in-medium width with the vacuum width in figure 4. Close to the resonance pole in-medium width and vacuum width are almost equal leading to an increase of the absorption cross section due to the contribution from the collision width. Because of the strong increase of the in-medium width with increasing Δ -mass the Δ -contribution is decreased at higher energies. This cancels the contribution from the collision width at photon energies of about 400 MeV .

In these calculations the Δ 's experience the same potential as the nucleons. If we apply a Δ -potential of $U_\Delta = -30 \rho/\rho_0 \text{ MeV}$ the Δ -peak is shifted to higher energies and decreased by about $70 \mu\text{b}$ per nucleon because the Δ -width increases strongly with increasing mass and the Δ -peak is proportional to $\frac{1}{\Gamma}$. Since the position of the Δ -peak becomes density dependent the integration over the volume of the nucleus leads to a further smearing out. It is important to note that a different parameterization of the Δ -width, for example a constant width, is not in line with the P_{33} -multipole of elastic πN -scattering [16]. Therefore the Δ -peak has to be reduced if shifted to higher

photon energies. Compared to the experimental data [2] we see from figure 7 that we underestimate the cross section in the high mass Δ -region.

The very different Δ -potential of $U_\Delta = 0$ MeV leads to a reduction of the Δ -peak by about $50 \mu\text{b}$ and thus does not give a better description of the experimental data.

Using a larger cross section for $\gamma n \rightarrow n \pi^0$ according to an alternative in [24] gives only a little improvement in the description of the experimental data.

In the region of the Δ -resonance a better description of the experimental cross section is only possible by an extension of our model. The missing strength indicates that we have to include further reaction mechanisms.

In the framework of the Δ -hole model there are calculations of the photoabsorption cross section up to energies of 500 MeV in [7] and up to 400 MeV in [8]. The results of both calculations differ by about 10%. The calculation in [7] gives a very good description of the experimental data. The authors claim that a two-body absorption process as depicted in figure 8 plays an important role in the energy range between 400 and 500 MeV. Our calculation of this diagram gives a similar result. Adding this contribution to the cross section described above leads to a good description of the experimental data (figure 9). Of course, just adding this one amplitude does not represent a gauge invariant calculation, but, in agreement with [7], it gives the right direction.

In the region of the higher resonances the effect of the in-medium widths is small since the collision widths are small as discussed in section 3. Here the collision widths essentially only compensate the Pauli reduction. At a photon energy of 700 MeV we still obtain a resonance-like structure and overestimate the experimental cross section by about 25%.

The main reason for the observed insensitivity of the $N(1520)$ -resonance region to the model ingredients lies thus in the fact that the collisional widths are rather small compared to the decay width. In a consistent treatment one also has to take into account the collision broadening of the nucleon ground state. This effect can be estimated from the imaginary parts of the selfenergies of nucleons in nuclear matter and amounts, depending on the nucleon momentum, to 0 - 60 MeV broadening with an average width being less than 10 MeV [25]. In order to mimic this effect and other possible in-

medium modifications we have used the ansatz:

$$\Gamma_{coll}^* = \Gamma_{coll} \left(\alpha + \beta \frac{\rho}{\rho_0} \right) \quad (24)$$

The Δ -peak is not only reduced with increasing α or β but also shifted to smaller photon energies since the collision width from $\Delta N \rightarrow N N$ does not - in contrast to the free width - vanish for small Δ -masses. For reasonable values of α and β the higher resonances are not affected.

In figure 7 (lower part) we also show the different contributions to the total cross section. Here we see that the rise of the cross section between 550 and 700 MeV is not only caused by the excitation of the D_{13} -resonance but also to a large extent by the opening of the two-pion background channel. The contribution of the one-pion channel shows almost no resonant structure in this energy regime.

Alberico et al. [9] and Kondratyuk et al. [10] have performed phenomenological fits to the photoabsorption cross section on nuclei in order to gain information about masses and widths of nucleon resonances in nuclei. The main result of these studies was the need for significantly increased resonance widths, up to a factor of 4 for the D_{13} -resonance to explain the observed absence of the higher lying resonances. If we use a collision width of about 300 MeV for the D_{13} then also in our calculation the resonant structure in the cross section disappears. However, as discussed in section 3, a collision width of 300 MeV for the D_{13} -resonance is hard to justify. Furthermore, in [9] the neglect of background terms as well as the assumption that the elementary cross section on the neutron is equal to the one on the proton has led these authors to an unsatisfactory determination of the resonance couplings to the $N\gamma$ -system in the vacuum. In [10], furthermore, the cross section of Armstrong et al. for $\gamma n \rightarrow X$ was used; as discussed in section 4 this cross section is probably incorrect. The fit in the region of the Δ -resonance gave a good result in [10] only because a constant Δ -width was used and Pauli blocking for the background was neglected. The fit of Kondratyuk et al. has been improved in a recent work by Bianchi et al. [3]. However, the 'background' as well as the neutron cross section are still treated in an unsatisfactory way.

While the collisional broadening can, according to our results, not explain the observed disappearance of the higher resonances, there are other in-medium effects that still have to be explored. For example, there is the possibility that the width of the D_{13} -resonance is increased by a strong medium

modification of the free width, for example caused by the strong coupling to the $N\rho$ -channel and a downward mass shift of the ρ -meson in the nuclear medium [26]. This may lead to the disappearance of the structure in the region of the D_{13} -resonance. Another possibility is a strong medium modification of the elementary $\gamma N \rightarrow N\pi\pi$ cross section in the nuclear medium or a medium effect on the background amplitudes.

An understanding of the disappearance of the D_{13} -resonance in the photoabsorption cross section might thus be possible by a comparison between theory and experiment with respect to more exclusive reaction channels [17]; for example, a measurement of two-pion photoproduction in nuclei would give information on the opening of the 2π -channel in the medium.

6 Summary and outlook

We have presented a calculation of the photoabsorption cross section on nuclei within a local Thomas-Fermi framework for photon energies from 300 MeV to 1 GeV. Starting from a reasonable parameterization of the free photon nucleon cross section we applied the medium modifications Fermi motion, Pauli blocking and collision broadening for the involved nucleon resonances.

The calculation of collision broadening required the knowledge of the interaction of nucleon resonances with nucleons. The cross section for $N R \rightarrow N N$ was obtained by detailed balance from one-pion production cross sections in nucleon-nucleon collisions under the assumption of a simple resonance model. Using this information we also estimated cross sections for $N R \rightarrow N R'$. For the $\Delta(1232)$ -resonance it turned out that collision broadening and reduction of the free width by Pauli blocking nearly compensate each other resulting in a very small net broadening. The collision broadening of the higher resonances in our model is almost negligible.

Our calculated photoabsorption cross section fails to describe the experimentally observed disappearance of the D_{13} -resonance. This might be caused by a strong broadening of the D_{13} -resonance in the nuclear medium due to a strong coupling to the $N\rho$ -channel and a mass shift of the ρ -meson in the nuclear medium but also by a medium modification of the $\gamma N \rightarrow N\pi\pi$ process.

An experimental measurement of exclusive cross sections would be helpful for a better understanding of the photon-nucleus reaction and an explana-

tion of the disappearance of the D_{13} -resonance in the total photonuclear absorption cross section. In particular, the experimental investigation of the 2π -channel on nuclei would be very important because of the opening of this channel in the $N(1520)$ -resonance region.

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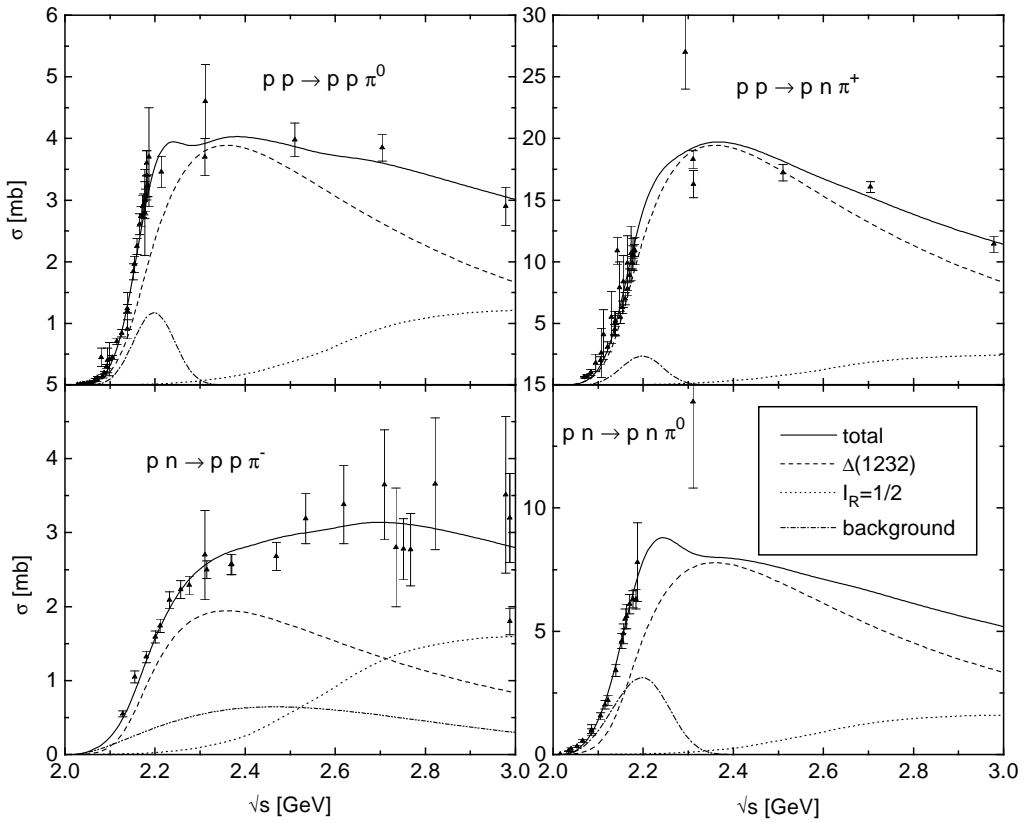


Figure 1: Resonance and background contributions to one-pion production cross sections in nucleon-nucleon collisions. The experimental data are taken from [21].

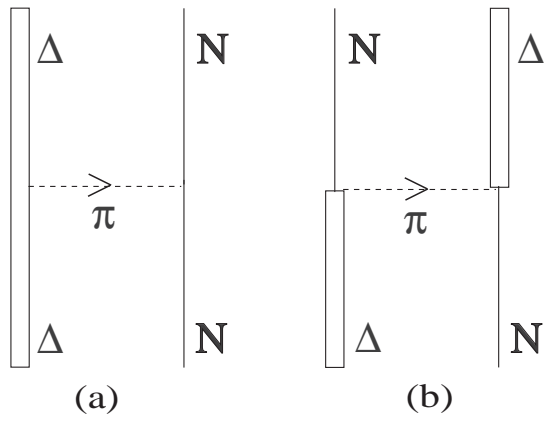


Figure 2: Diagrams to $\Delta N \rightarrow \Delta N$.

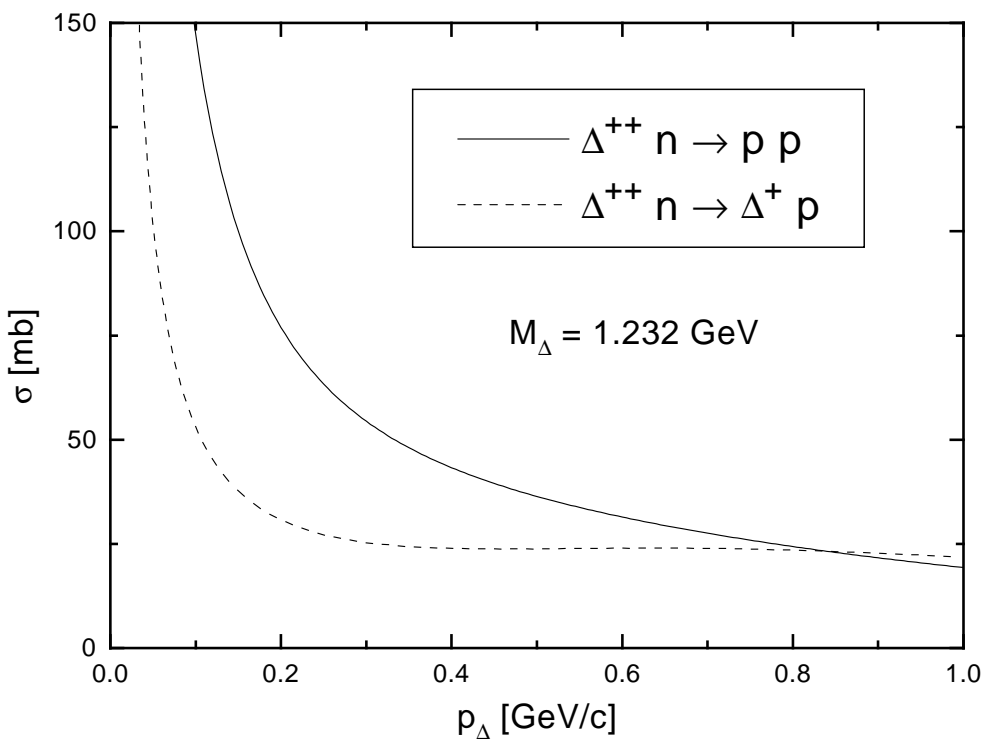


Figure 3: Cross sections for Δ -nucleon collisions. p_Δ is the Δ -momentum in the rest frame of the nucleon.

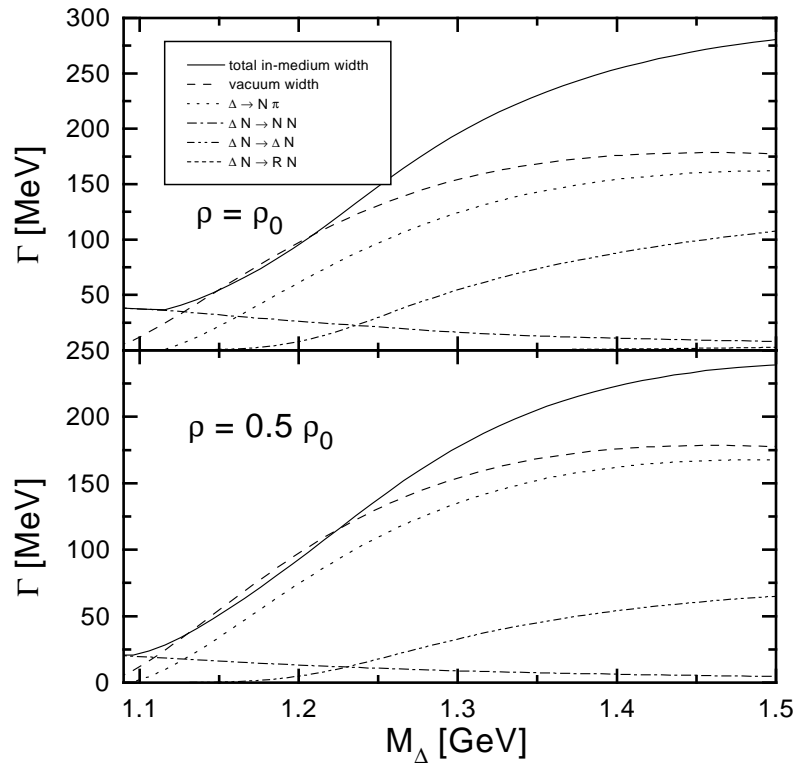


Figure 4: In-medium width of the $\Delta(1232)$.

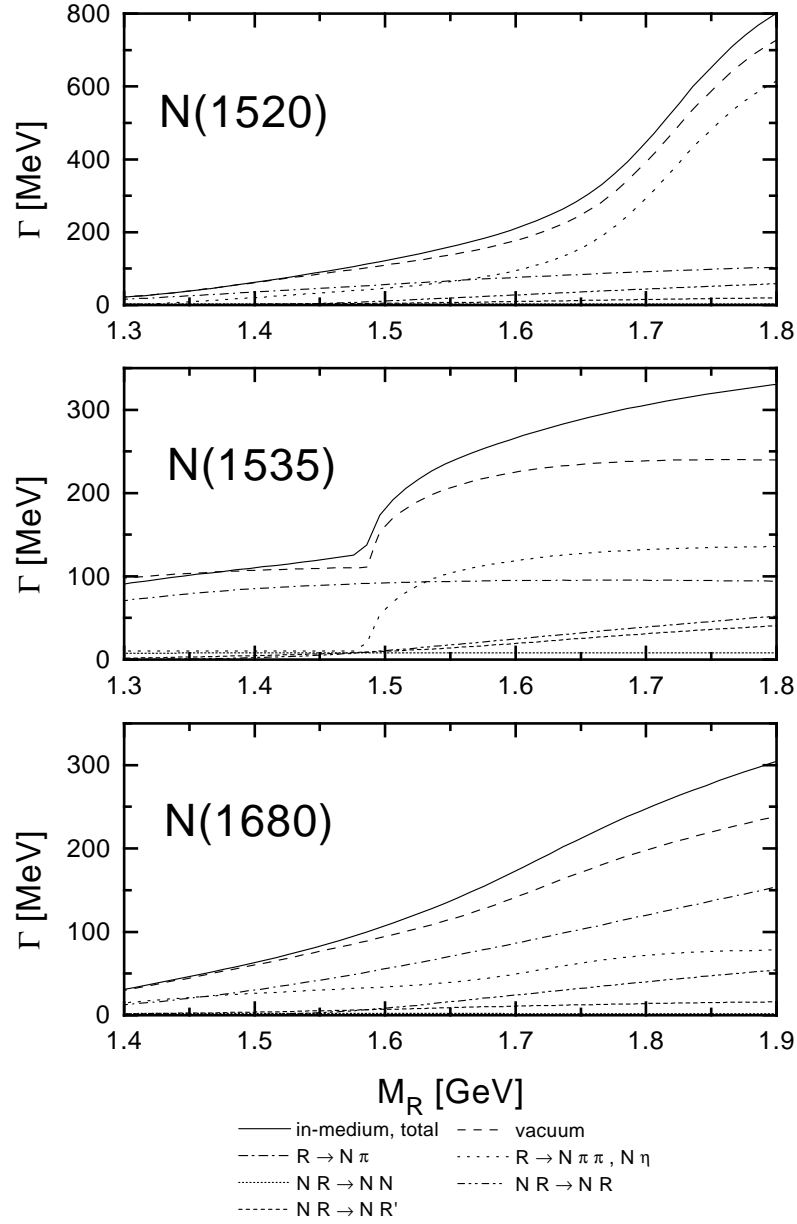


Figure 5: In-medium widths of the higher resonances that are relevant for photonuclear processes.

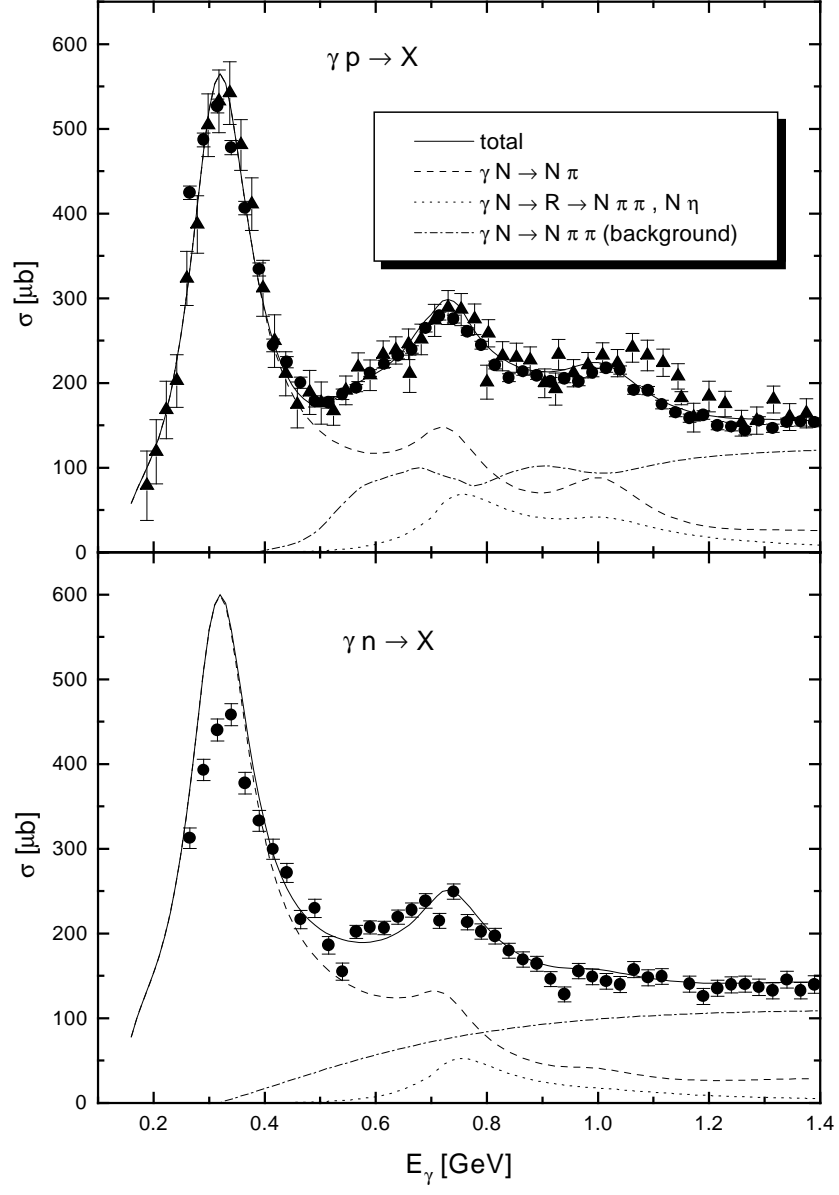


Figure 6: Total photoabsorption cross sections on the nucleon. The experimental proton data are from [5] (circles) and [21] (triangles). The neutron data are taken from [6].

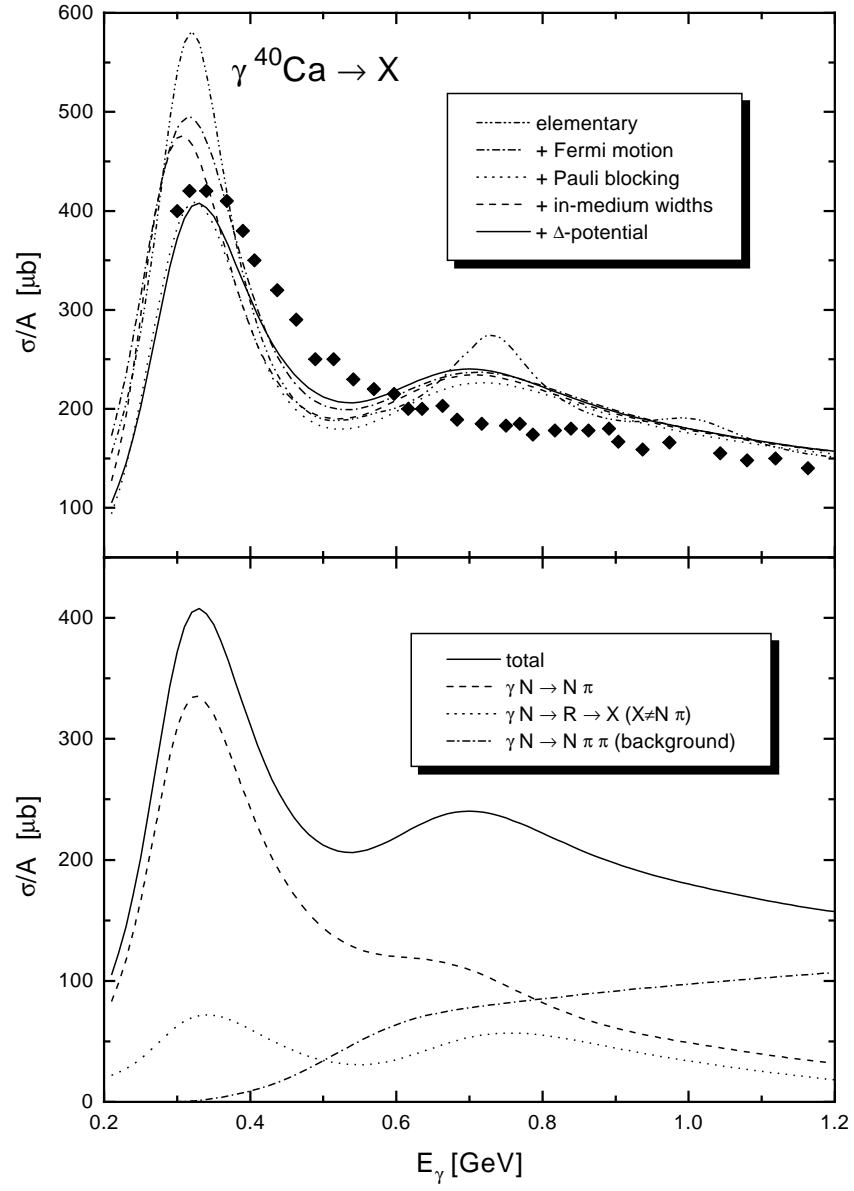


Figure 7: Photoabsorption cross section on ${}^{40}\text{Ca}$. The experimental data are obtained by an average over different nuclei [2].

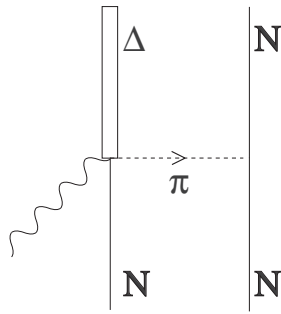


Figure 8: Example of a diagram contributing to $\gamma N N \rightarrow N \Delta(1232)$.

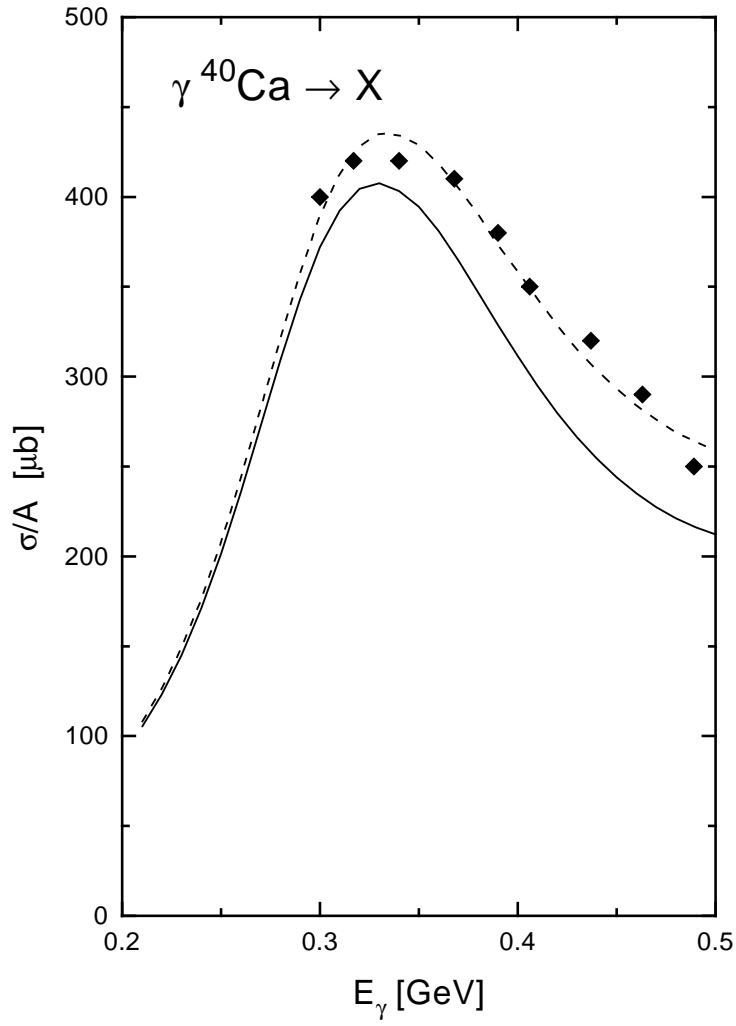


Figure 9: Total photoabsorption cross section on ^{40}Ca with the additional two-body absorption mechanism $\gamma N N \rightarrow N \Delta(1232)$ of figure 8 (dashed line). The solid line represents the previous cross section.